

# Natural emergence of neutrino masses and bi-linear “RPV” from $R$ -symmetry

Sabyasachi Chakraborty<sup>a</sup>, Joydeep Chakraborty<sup>b</sup>

<sup>a</sup>*Department of Theoretical Physics, Tata Institute of Fundamental Research, Mumbai 400005, India*

<sup>b</sup>*Department of Physics, Indian Institute of Technology, Kanpur-208016, India*

---

## Abstract

We propose a supersymmetric extension of the Standard Model (SM) with a continuous global  $U(1)_R$  symmetry. The  $R$ -charges of the SM fields are identified with that of their lepton numbers. As an artifact, a bi-linear “ $R$ -parity violating” term emerges at the superpotential. However,  $R$ -symmetry is not an exact symmetry as it is broken by supergravity effects. As a result, sneutrinos acquire a small vacuum expectation value in this framework. Consequently, neutrino masses and mixing can be explained at the tree level itself. This is noticeably different from the standard  $R$ -parity violating Minimal Supersymmetric Standard Model. Gravitino mass turns out to be the order parameter of  $R$ -breaking and directly related to the neutrino mass. To have thermal production of gravitinos, we require the gravitino mass to be  $> 200$  keV. We show that such a gravitino is an excellent dark matter candidate and is safe from all constraints. Finally, we looked into the collider implications of our framework.

---

## 1. Introduction

Neutrino oscillation experiments have firmly established the existence of tiny non-zero masses of active neutrinos and non-trivial mixing in the lepton sector. Neutrinos are massless within the paradigm of Standard Model (SM). Therefore, the oscillation data can naturally promote the query for physics beyond the SM (BSM). Apart from CP-violating phases most of the oscillation parameters have been determined with great accuracy [1]. Thus, it is expected that any BSM theory must provide a suitable mechanism for neutrino mass generation. To do so, see-saw mechanism seems to be the natural one that predicts the Majorana nature of the neutrino signifying lepton number violation. The basic idea behind the see-saw mechanism is to integrate out the heavy modes, leading to higher-dimensional neutrino mass operators. Depending on the choice of the heavy particles, one can classify variants of this mechanism, namely Type-I, -II, -III, Inverse, Double see-saw *etc.* both in supersymmetric (SUSY) and non-SUSY scenarios. However, so far there is no such experimental data that can help us to pin down any particular mechanism belonging to any specific model. Thus from neutrino mass generation perspective, these mechanisms are more sacred than particular models. Apart from neutrino masses and mixing the deviation from the galactic rotation curve and bullet clusters provide concrete evidences in favour of dark matter (DM) which remains unexplained within the SM framework. Cosmological observations has also measured the relic density of the DM with very high precision [2, 3]. But unfortunately the DM characteristics, e.g. mass, spin, and its nature, i.e., cold, warm, single or multi component are yet to be determined. Thus there must be a theory beyond the Standard Model.

Keeping these things in mind and also from a theoretical perspective, SUSY demands special attention as a potential BSM theory. Apart from prescribing a solution for the hierarchy problem which SM suffers from,

---

*Email addresses:* [sabya@theory.tifr.res.in](mailto:sabya@theory.tifr.res.in) (Sabyasachi Chakraborty), [joydeep@iitk.ac.in](mailto:joydeep@iitk.ac.in) (Joydeep Chakraborty)

it can also explain neutrino masses-mixing and provide a suitable DM candidate. For example, the lightest supersymmetric particle (LSP) is an excellent DM candidate in the paradigm of the minimal supersymmetric standard model (MSSM) with  $R$ -parity conservation. On the other hand MSSM with  $R$ -parity violation (RPV) is an intrinsic SUSY way to generate neutrino masses and mixing both at the tree level as well as at the one-loop level [4]. However, if  $R$ -parity is broken, the LSP becomes unstable and hence fails to explain the observed relic density of the Universe. But there could be potential dark matter candidates, in the form of a gravitino, axino, axion and keV sterile neutrino [5].

Concerning the experimental verification of proposed models, unfortunately, the early 13 TeV run of the LHC has not found any signals in favour of SUSY [6, 7]. Although the non-observation of superpartners does not invalidate the idea of SUSY, but it certainly questions the ability of MSSM to resolve the naturalness problem. For example, LHC has already ruled out gluinos lighter than 2 TeV when the gluino and LSP masses are well separated. This in turn makes the whole coloured sector heavy due to the logarithmic sensitivity to the ultra-violet (UV) scale through renormalisation group evolution equations. Interestingly this correlation is not generic and can be avoided within the models with  $R$ -symmetry and Dirac gauginos. One needs to extend the gauge sector of MSSM to  $N = 2$  representation to construct a Dirac gaugino mass. This requires a singlet  $\hat{S}$ , a triplet  $\hat{T}$  under  $SU(2)_L$  and an octet  $\hat{O}$  under  $SU(3)_C$  living in the adjoint representation of the SM gauge group. These fields couple with bino, wino and gluinos respectively, to generate Dirac masses for the gauginos. The presence of additional adjoint scalars cancel the UV logarithmic divergence, resulting in a finite correction [8] only. Hence, the Dirac gluino masses can be easily made heavy. An immediate consequence of having heavy Dirac gluinos is the suppression of the cross-section for gluino pair or squark-gluino associated pair productions. In addition, the pair production of squarks are also suppressed as it requires chirality flipping Majorana gaugino masses in the propagator which these frameworks are devoid of. This invariably weakens [9] the stringent upper bound on the first two generation squark masses. The added feature of  $R$ -symmetry, apart from excluding Majorana masses for gauginos, is to prohibit trilinear scalar couplings ( $A$ -terms) and the traditional higgsino mass term, i.e., the  $\mu$  term. The absence of  $A$ -terms result in the suppression of the flavor and CP violating interactions [10].

Motivated by these nice features, we propose a SUSY scenario enveloped by  $R$ -symmetry. Our prime aim in this paper is to generate neutrino masses within the  $R$ -symmetric Dirac gaugino framework. In fact the  $R$ -charges are now identified with the lepton number of the SM fields. In general, we have the freedom to cast these charges and one can look into other assignments in [11–19]. An artifact of our  $R$ -charge assignment is the presence of a bi-linear “ $R$ -parity violating” term in the superpotential which is indeed  $R$ -symmetric. This is a very striking feature as it allows the sneutrino to acquire vacuum expectation value (VEV) when  $R$ -symmetry gets broken. This leads to a mixing between light neutrinos and Dirac gauginos which in turn helps to generate light neutrino masses at the tree-level. This RPV coupling will allow the gravitino to decay to neutrino and photons. Such a gravitino is an excellent decaying dark matter candidate provided its lifetime is greater than the age of the universe.

We categorize our paper as follows. To start with, we propose the model, and discuss the basic features of this scenario. We emphasize on the specific choice of  $R$ -charges which lead to the emergence of bi-linear “RPV” terms in the superpotential, which are  $R$ -symmetric. Then in the following section, sec. (2.2), we discuss soft SUSY breaking,  $R$ -symmetry breaking terms, and also the generation of sneutrino VEV. In the next section we explore the neutral fermion mass matrix. There we show, how successfully the neutrino masses and mixing can be generated in a simplistic scenario. We discuss the more generic methods in appendix (7). While considering normal as well as inverted hierarchies, we note the constraints on the relevant superpotential parameters. As it happens, the gravitino is the order parameter of  $R$ -symmetry breaking. We also explore the possibility of gravitino dark matter in our model satisfying necessary constraints. Altogether, we come up with a conclusion that after successfully defending relevant constraints, our model can describe the neutrino mass and mixing parameters. We conclude by providing a direction to adjudicate this scenario in collider experiments.

## 2. The Model

Our scenario is based on a SUSY framework with an added global  $R$ -symmetry where the gauginos are Dirac type unlike the MSSM. Enjoying the freedom of  $R$ -charge assignments, we tag the chiral superfields with very specific values as depicted in table 1. It is rather straightforward to show that the scalars share the

Superfields	SM rep	$U(1)_R$	Superfields	SM rep	$U(1)_R$
$\widehat{H}_u$	$(1, 2, 1)$	0	$\widehat{R}_d$	$(1, 2, -1)$	2
$\widehat{H}_d$	$(1, 2, -1)$	0	$\widehat{R}_u$	$(1, 2, 1)$	2
$\widehat{Q}_i$	$(3, 2, \frac{1}{3})$	1	$\widehat{S}$	$(1, 1, 0)$	0
$\widehat{U}_i^c$	$(\bar{3}, 1, -\frac{4}{3})$	1	$\widehat{T}$	$(1, 3, 0)$	0
$\widehat{D}_i^c$	$(\bar{3}, 1, \frac{2}{3})$	1	$\widehat{O}$	$(8, 1, 0)$	0
$\widehat{L}_i$	$(1, 2, -1)$	2			
$\widehat{E}_i^c$	$(1, 1, 2)$	0			

Table 1: The gauge quantum numbers under the SM gauge group  $SU(3)_c \times SU(2)_L \times U(1)_Y$  as well as the  $U(1)_R$  charge assignments of the chiral superfields residing in the model.

same  $R$ -charges with their corresponding superfields whereas for fermions they are one less. Similarly, the gauginos have  $R$ -charge one and the corresponding gauge bosons have zero  $R$ -charge. The lepton number of the SM particles are identified with their  $R$ -charges as shown in table 1. The lepton number of the superpartners can be non-standard. The definition of our  $R$ -charge assignment can also be understood as, our choice of  $R$  equals  $R_0 + L$  [13] where  $R_0$  is the traditional  $R$ -charge assignment in MRSSM [20–23]<sup>1</sup> and  $L$  stands for the lepton number. As we know, an invariant action demands the superpotential to have  $R$ -charge of two units. Hence, the allowed terms in superpotential are:

$$\begin{aligned}
W_{\text{MSSM}} &= y_u^{ij} \widehat{H}_u \widehat{Q}_i \widehat{U}_j^c + y_d^{ij} \widehat{H}_d \widehat{Q}_i \widehat{D}_j^c + y_e^{ij} \widehat{L}_i \widehat{H}_d \widehat{E}_j^c + y_r^i \widehat{R}_d \widehat{H}_d \widehat{E}_i^c, \\
W_{\text{adj}} &= \lambda_S^i \widehat{S} \widehat{H}_u \widehat{R}_d + \lambda_T^u \widehat{H}_u \widehat{T} \widehat{R}_d + \lambda_S^d \widehat{S} \widehat{H}_d \widehat{R}_u + \lambda_T^d \widehat{H}_d \widehat{T} \widehat{R}_u + \xi_i \widehat{S} \widehat{H}_u \widehat{L}_i + \eta_i \widehat{H}_u \widehat{T} \widehat{L}_i, \\
W_{\text{“bi-RPV”}} &= \epsilon_i \widehat{H}_u \widehat{L}_i, \\
W_\mu &= \mu_u \widehat{H}_u \widehat{R}_d + \mu_d \widehat{H}_d \widehat{R}_u.
\end{aligned} \tag{1}$$

The total superpotential of our scenario is  $W = W_{\text{MSSM}} + W_{\text{adj}} + W_{\text{“bi-RPV”}} + W_\mu$ . The triplet  $\widehat{T}$  under  $SU(2)_L$  is parametrised as  $\widehat{T} = \sum_{a=1,2,3} \widehat{T}^{(a)}$ , where  $\widehat{T}^{(a)} = T_a \sigma^a / 2$ ,  $\sigma^a$ 's being the Pauli matrices and we also denote the components of the triplet field as  $T_3 = T_0$ ,  $T_+ = (T_1 - iT_2)/\sqrt{2}$  and  $T_- = (T_1 + iT_2)/\sqrt{2}$ .  $\widehat{H}_u, \widehat{H}_d, \widehat{L}_i, \widehat{E}_i^c, \widehat{Q}_i, \widehat{U}_i^c, \widehat{D}_i^c$  are the prototypical MSSM fields.  $\lambda^u, \lambda^d, \xi_i, \eta_i, y_u, y_d, y_e$  are trilinear/Yukawa couplings and  $\mu, \epsilon$  are couplings with mass dimension. Notice that, the lack of a  $\mu \widehat{H}_u \widehat{H}_d$  term is a serious problem as the higgsino like chargino masses directly stem from this. The LEP experiment has ruled out charginos below 100 GeV. Hence, to generate a  $\mu$  term, it is mandatory to include two additional  $SU(2)$ -doublet chiral superfields  $\widehat{R}_u$  and  $\widehat{R}_d$  carrying non-zero  $R$ -charges. The presence of a  $R$ -charge for these two fields imply that  $R$ -symmetry cannot be broken in the visible sector spontaneously otherwise we have to encounter massless  $R$ -axions. As a result, these doublets are known as inert doublets in the literature.

We like to draw attention towards  $W_{\text{“bi-RPV”}}$  which contains the bi-linear “RPV” but  $R$ -symmetric terms. These terms are automatically appearing as a relic of our  $R$ -charge assignment. It is worthy to mention that in our model,  $R$ -symmetry prohibits baryonic “RPV” terms in the superpotential and in the process the stringent constraints from proton decay are taken care of. Before discussing the neutrino mass generation we would like to first address soft SUSY breaking and  $R$ -symmetry breaking leading to the generation of sneutrino VEV.

<sup>1</sup>We note in passing that  $R_0$  is one for the lepton and quark superfields.

### 2.1. Soft (super-soft) SUSY breaking interactions

We choose to work in a scenario where SUSY (global) breaking is not associated with  $R$ -symmetry breaking. This can be achieved through both  $D$ - and  $F$ -type spurions. For example, Dirac gaugino masses can be generated with the help of a spurion superfield  $W'_\alpha = \lambda'_\alpha + \theta_\alpha D'$  which appears in the Lagrangian as follows [24, 25]:

$$\mathcal{L}_{\text{Dirac}}^{\text{gaugino}} = \int d^2\theta \frac{W'_\alpha}{\Lambda} \left[ \kappa_1 W_{1\alpha} \hat{S} + \kappa_2 \text{Tr}(W_{2\alpha} \hat{T}) + \kappa_3 \text{Tr}(W_{3\alpha} \hat{O}) \right] + h.c., \quad (2)$$

where  $W_{i\alpha}$  and  $W'_{i\alpha}$  have  $R$ -charge 1, i.e.,  $R[\lambda_{i\alpha}] = R[\lambda'_{i\alpha}] = 1$  and  $R[D'] = 0$ . The integration over the Grassmann co-ordinates generate the Dirac gaugino masses  $M_i^D \sim \kappa_i \langle D' \rangle / \Lambda$ . Here, the  $i = 1, 2, 3$  for  $U(1)_Y$ ,  $SU(2)_L$  and  $SU(3)_C$  respectively, and  $\Lambda$  represents the messenger scale. After the discovery of Higgs boson around 125 GeV it is important to address the status of Higgs mass within the given scenario along with other scalar masses as well. In a purely supersoft scenario, the scalar masses are one-loop suppressed as opposed to the gaugino masses. As an outcome,  $D$ -flatness is a natural direction [8] which leads to a vanishingly small tree level quartic for the Higgs field. This is problematic from the perspective of fitting the Higgs mass to 125 GeV. Thus instead of working in the generalised supersoft supersymmetry framework [26], we recall  $F$ -type breaking [13, 15]. In our model, the scalar masses can be generated through a  $F$ -type spurion defined as  $\hat{X} = x + \theta^2 F_X$  which allows the following  $U(1)_R$  preserving operators [13]

$$\begin{aligned} & \int d^4\theta \frac{\hat{X}^\dagger \hat{X}}{\Lambda^2} \left[ \sum_i \hat{\Phi}_i^\dagger \hat{\Phi}_i + \left\{ \hat{H}_u \hat{H}_d + \epsilon \Lambda \hat{S} + \hat{S}^2 + \hat{T}^2 + \left( \frac{1}{\Lambda} \times \text{cubic} \right) + h.c. \right\} \right], \\ & \int d^2\theta \frac{X}{\Lambda} \left( \hat{S} \hat{T}^2 + \hat{S} \hat{O}^2 + \hat{S}^3 \right) + h.c., \end{aligned} \quad (3)$$

and can automatically generate the following  $U(1)_R$  preserving renormalizable soft SUSY breaking terms

$$\mathcal{L}_{\text{soft}} = \sum_i m_i^2 \phi_i^\dagger \phi_i + \left[ t_S S + \frac{1}{2} b_S S^2 + B_\mu H_u H_d + \dots \right]. \quad (4)$$

The soft mass squared terms are proportional to  $\langle F_X \rangle^2 / \Lambda^2 \equiv M_{\text{SUSY}}^2$  and we also consider same magnitude for the  $F$ - and  $D$ -type VEVs. In passing, we would like to note that such a mechanism also generates a scalar singlet tadpole  $t_S S$ . However, as long as  $t_S < M_{\text{SUSY}}^2$ , such a tadpole is not expected to destabilize the hierarchy [27]. Nevertheless, due to the absence of  $R$ -breaking terms  $B_\epsilon H_u \tilde{\ell}_i$  in the scalar sector, sneutrinos cannot acquire VEV which is a relevant ingredient for neutrino mass generation. This has been a requirement for us to consider the possibility of  $R$ -symmetry breaking.

### 2.2. $R$ -symmetry breaking

It is well established that our Universe is associated with a vanishingly small vacuum energy or cosmological constant. To explain this from the perspective of spontaneously broken supergravity theory in the hidden sector, the superpotential needs to acquire a non-zero VEV. Since the superpotential carries non-zero  $R$ -charge, therefore,  $\langle W \rangle \neq 0$  implies breaking of  $R$ -symmetry. As a result, the gravitino would also acquire a mass. Here, the order parameter of  $R$ -breaking is the gravitino mass. This  $R$ -breaking information is communicated to the visible sector through anomaly mediation and in the process the following  $R$ -symmetry breaking terms are generated

$$\begin{aligned} \mathcal{L}_{\mathcal{R}} &= M_1 \tilde{b} \tilde{b} + M_2 \tilde{w} \tilde{w} + M_3 \tilde{g} \tilde{g} \\ &+ A_u \tilde{u}_R \tilde{u}_L^* H_u^0 + A_d \tilde{d}_R \tilde{d}_L^* H_d^0 + A_\ell \tilde{\ell}_R \tilde{\ell}_L^* H_d^0 + h.c., \end{aligned} \quad (5)$$

where the Majorana gaugino masses are generated through small  $R$ -breaking effects as

$$M_i = \frac{g_i^2}{16\pi^2} b_i m_{3/2} \quad (i = 1, 2, 3), \quad (6)$$

with beta functions

$$b_1 = 33/5, \quad b_2 = 1, \quad b_3 = -3. \quad (7)$$

The small  $R$  symmetry-breaking trilinear scalar interactions are as follows

$$\begin{aligned} A_\tau &= \frac{m_{3/2}}{16\pi^2} \left( -\frac{9}{5}g_1^2 - 3g_2^2 + 3Y_b^2 + 4Y_\tau^2 \right), \\ A_t &= \frac{m_{3/2}}{16\pi^2} \left( -\frac{13}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + 6Y_t^2 + Y_b^2 \right), \\ A_b &= \frac{m_{3/2}}{16\pi^2} \left( -\frac{7}{15}g_1^2 - 3g_2^2 - \frac{16}{3}g_3^2 + Y_t^2 + 6Y_b^2 + Y_\tau^2 \right). \end{aligned} \quad (8)$$

It is also important to note that the natural presence of a conformal compensator field,  $\Sigma = 1 + \theta^2 m_{3/2}$ , invariably generates a  $B\epsilon_i$  term in the superpotential through the following operator [28]:

$$\mathcal{L} = \int d^2\theta \Sigma^3 \epsilon_i \hat{H}_u \hat{L}_i. \quad (9)$$

After scaling out this compensator field with  $\hat{\Phi}' = \Sigma \hat{\Phi}$ , where  $\hat{\Phi}$  is a chiral superfield, we generate a bi-linear term  $(H_u \tilde{\ell}_i)$  in the scalar potential

$$B\epsilon_i = \epsilon_i m_{3/2}. \quad (10)$$

Hence, the  $B\epsilon$  term is always aligned with the  $\epsilon$  term. Such terms are  $R$ -breaking effects and proportional to the gravitino mass. The presence of this small effect would invariably generate a tiny sneutrino VEV which is appropriate for neutrino mass-mixing, as we discuss in the next section.

### 2.3. Sneutrino VEV

To compute the sneutrino VEV, one has to include the contributions from  $F$ -,  $D$ - and soft SUSY breaking terms. The additional pieces associated with  $SU(2)_L$  and  $U(1)_Y$  in the  $D$ -terms are

$$D_2^a = g \left( H_u^\dagger \tau^a H_u + \tilde{\ell}_i^\dagger \tau^a \tilde{\ell}_i + T^\dagger \lambda^a T \right) + \sqrt{2} (M_2^D T^a + M_2^D T^{a\dagger}), \quad (11)$$

where  $\tau^a$  and  $\lambda^a$ 's represent the  $SU(2)$  generators in the fundamental and adjoint representations respectively. Similarly, the weak hyper-charge contribution  $D_Y$  is given by

$$D_Y = \frac{g'}{2} \left( H_u^\dagger H_u - \tilde{\ell}_i^\dagger \tilde{\ell}_i \right) + \sqrt{2} M_1^D (S + S^\dagger). \quad (12)$$

The tree level potential terms which participate in the sneutrino field minimisation equations are

$$\begin{aligned} V_F &= |\epsilon_i|^2 |\tilde{\nu}_i^0|^2, \\ V_{\text{soft}} &= \tilde{m}^2 |\tilde{\nu}_i^0|^2 + B\epsilon_i H_u^0 \tilde{\nu}_i, \\ V_D &= \left[ \frac{(g^2 + g'^2)}{8} |\tilde{\nu}_i^0|^2 - \sqrt{2} g' M_1^D v_S + \sqrt{2} g M_2^D v_T \right] |\tilde{\nu}_i^0|^2. \end{aligned} \quad (13)$$

In the limit of vanishing VEV of the singlet ( $S$ ) and the triplet ( $T$ ), i.e.,  $v_S, v_T \rightarrow 0$ , the sneutrino VEV can be approximated as

$$\langle \tilde{\nu}_i \rangle = -\frac{B\epsilon_i v_u}{\tilde{m}_i^2 + \epsilon_i^2}. \quad (14)$$

Such a choice of the singlet and the triplet VEVs also ensure that these fields are very heavy. Assuming  $\langle H_u^0 \rangle = v_u \sim \tilde{m}_i$  i.e., at the electroweak scale, we find  $\langle \tilde{\nu}_i \rangle \sim B\epsilon_i / \tilde{m} \sim \epsilon_i m_{3/2} / \tilde{m} \equiv \tilde{\epsilon}_i m_{3/2}$ <sup>2</sup>. In the next section, we investigate the neutral fermion sector which is of main importance of this analysis.

<sup>2</sup>Off course, in the same manner the inert scalars ( $R_u, R_d$ ) would also acquire a VEV due to  $R$ -symmetry breaking and as a result would mix with the Higgs fields. However, that mixing is also suppressed by the  $R$ -breaking parameter gravitino mass and does not play any important role in the phenomenological description.

### 3. The neutral fermion sector

The Lagrangian corresponding to the neutral fermion sector after  $R$ -symmetry breaking contain the following terms

$$\begin{aligned}
\mathcal{L}_{f^0} = & M_1^D \tilde{b} \tilde{S} + M_2^D \tilde{w}^0 \tilde{T} + \mu_u \tilde{H}_u^0 \tilde{R}_d^0 + \mu_d \tilde{H}_d^0 \tilde{R}_u^0 + \lambda_S^u v_u \tilde{S} \tilde{R}_d^0 - \lambda_S^d v_d \tilde{S} \tilde{R}_u^0 + \lambda_T^u v_u \tilde{T}^0 \tilde{R}_d^0 + \lambda_T^d v_d \tilde{T}^0 \tilde{R}_u^0 \\
& + M_1 \tilde{b} \tilde{b} + M_2 \tilde{w}^0 \tilde{w}^0 + M_S \tilde{S} \tilde{S} + M_T \tilde{T}^0 \tilde{T}^0 + \frac{g' v_u}{\sqrt{2}} \tilde{b} \tilde{H}_u^0 - \frac{g' v_d}{\sqrt{2}} \tilde{b} \tilde{H}_d^0 - \frac{g v_u}{\sqrt{2}} \tilde{w}^0 \tilde{H}_u^0 + \frac{g v_d}{\sqrt{2}} \tilde{w}^0 \tilde{H}_d^0 \\
& - \frac{g' v_i}{\sqrt{2}} \tilde{b} \tilde{\nu}_i + \xi_i v_u \tilde{S} \nu_i + \frac{g v_i}{\sqrt{2}} \tilde{w}^0 \tilde{\nu}_i + \eta_i v_u \tilde{T}^0 \nu_i + (\epsilon_i + \xi_i v_S + \eta_i v_T) \tilde{H}_u^0 \nu_i.
\end{aligned} \tag{15}$$

Here,  $i$  stands for  $e, \mu$  and  $\tau$  respectively.

The neutral fermion mass matrix in the basis  $f^0 = (\nu_i, \tilde{b}, \tilde{S}, \tilde{w}^0, \tilde{T}^0, \tilde{H}_u^0, \tilde{R}_d^0, \tilde{H}_d^0, \tilde{R}_u^0)$  has the schematic form

$$\mathcal{L}_{f^0}^{\text{mass}} = \frac{1}{2} (f^0)^T M_N f^0. \tag{16}$$

where

$$M_N = \begin{pmatrix} m_{\tilde{f}^0} & m_D \\ m_T^D & 0 \end{pmatrix}, \tag{17}$$

with

$$m_{\tilde{f}^0} = \begin{pmatrix} M_1 & M_1^D & 0 & 0 & \frac{g' v_u}{\sqrt{2}} & 0 & -\frac{g' v_d}{\sqrt{2}} & 0 \\ M_1^D & M_S & 0 & 0 & \xi_i v_i & \lambda_S^u v_u & 0 & \lambda_S^d v_d \\ 0 & 0 & M_2 & M_2^D & -\frac{g v_u}{\sqrt{2}} & 0 & \frac{g v_d}{\sqrt{2}} & 0 \\ 0 & 0 & M_2^D & M_T & \eta_i v_i & \lambda_T^u v_u & 0 & \lambda_T^d v_d \\ \frac{g' v_u}{\sqrt{2}} & \xi_i v_i & -\frac{g v_u}{\sqrt{2}} & \eta_i v_i & 0 & \mu & 0 & 0 \\ 0 & \lambda_S^u v_u & 0 & \lambda_T^u v_u & \mu & 0 & 0 & 0 \\ -\frac{g' v_d}{\sqrt{2}} & 0 & \frac{g v_d}{\sqrt{2}} & 0 & 0 & 0 & 0 & \mu \\ 0 & -\lambda_S^d v_d & 0 & \lambda_T^d v_d & 0 & 0 & \mu & 0 \end{pmatrix}, \tag{18}$$

and

$$m_D = \begin{pmatrix} -\frac{g' v_1}{\sqrt{2}} & -\frac{g' v_2}{\sqrt{2}} & -\frac{g' v_3}{\sqrt{2}} \\ \xi_1 v_u & \xi_2 v_u & \xi_3 v_u \\ \frac{g v_1}{\sqrt{2}} & \frac{g v_2}{\sqrt{2}} & \frac{g v_3}{\sqrt{2}} \\ \eta_1 v_u & \eta_2 v_u & \eta_3 v_u \\ (\epsilon_1 + \xi_1 v_S + \eta_1 v_T) & (\epsilon_2 + \xi_2 v_S + \eta_2 v_T) & (\epsilon_3 + \xi_3 v_S + \eta_3 v_T) \\ 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}. \tag{19}$$

Here,  $\langle H_u^0 \rangle = v_u$  and  $\langle H_d^0 \rangle = v_d$ , with the constraint that  $v = \sqrt{v_u^2 + v_d^2}$  and  $\tan \beta = v_u/v_d$ . We also consider  $\mu_u \equiv \mu_d \equiv \mu$ , and  $M_1 \sim M_2 \sim M_S \sim M_T \sim m_{3/2}/16\pi^2$ . We note that an order one value of the superpotential coupling  $\lambda_{S,T}^u$  provides substantial one-loop correction to the up-type Higgs boson mass  $m_{H_u}^2$  [29, 30]. These corrections are ‘stop-like’ and an 125 GeV Higgs boson can be obtained without requiring too heavy top squarks. In the next section, to carry on our proposed analysis, we adopt a simplified framework after integrating out some suitably chosen heavy fields. This helps us to understand the phenomenological aspects in detail.

### 3.1. Neutrino mass and mixing

In the limit when both the wino and the higgsinos are heavy and effectively decoupled, the neutral fermion mass matrix can be written in the basis  $(\nu_i, \tilde{b}, \tilde{S})$ <sup>3</sup>. We neglect the bino and singlino Majorana masses which are proportional to  $m_{3/2}/16\pi^2$  as their inclusion does not alter our conclusion. The neutral fermion mass matrix in such a limit can be depicted as

$$M_N = \begin{pmatrix} 0 & -\frac{g'}{\sqrt{2}}\tilde{\epsilon}_i m_{3/2} & \xi_i v_u \\ -\frac{g'}{\sqrt{2}}\tilde{\epsilon}_i m_{3/2} & 0 & M_D \\ \xi_i v_u & M_D & 0 \end{pmatrix}. \quad (20)$$

In order to diagonalise this effective matrix we adopt the methodology of [31] and choose  $\xi_i = \xi \mathbf{u}$ ,  $\tilde{\epsilon}_i = \tilde{\epsilon} \mathbf{v}$ . This results in

$$\begin{aligned} \xi_i &= \frac{\xi e^{i\theta/2}}{\sqrt{2}} \left[ \sqrt{1+\rho} U_{i3(i2)} + \sqrt{1-\rho} U_{i2(i1)} \right], \\ \tilde{\epsilon}_i &= \frac{\tilde{\epsilon} e^{-i\theta/2}}{\sqrt{2}} \left[ \sqrt{1+\rho} U_{i3(i2)}^* - \sqrt{1-\rho} U_{i2(i1)}^* \right], \end{aligned} \quad (21)$$

where  $\rho = \frac{\sqrt{1+r}-\sqrt{r}}{\sqrt{1+r}+\sqrt{r}}$  with  $r \equiv \frac{|\Delta m_{12}|^2}{|\Delta m_{23}|^2}$ . The indices (in the parenthesis) denote elements of the PMNS matrix for normal (inverted) hierarchy. Using this approach we find out the suitable parameters for both cases.

#### 3.1.1. Normal hierarchy

For simplicity, we choose the lightest neutrino mass ( $m_1$ ) to be zero. The other mass eigenvalues can be readily obtained as  $m_2 = g'\tilde{\epsilon}m_{3/2}\xi(1-\rho)/\sqrt{2}$  and  $m_3 = g'\tilde{\epsilon}m_{3/2}\xi(1+\rho)/\sqrt{2}$ . With the help of present oscillation parameters [1] we find

$$\tilde{\epsilon} \xi m_{3/2} \sim 10^{-10} \text{ GeV}. \quad (22)$$

Using eq. (21), the parameters should have the following structures to be consistent with the neutrino oscillation data

$$\xi_i = \frac{\xi}{\sqrt{2}} \begin{pmatrix} 0.487 \\ 1.165 \\ 0.635 \end{pmatrix}, \quad \tilde{\epsilon}_i = \frac{\tilde{\epsilon}}{\sqrt{2}} \begin{pmatrix} -0.1 \\ 0.55 \\ 1.29 \end{pmatrix}. \quad (23)$$

For simplicity we assume the CP violating Dirac phase to be zero and all the  $R$ -conserving masses ( $M_D, v_u$ ) to be at the electroweak scale.

#### 3.1.2. Inverted hierarchy

In case of inverted hierarchy we follow the same prescription and find  $m_3 = 0$ ,  $m_2 = g'\tilde{\epsilon}m_{3/2}\xi(1+\rho)/\sqrt{2}$  and  $m_1 = g'\tilde{\epsilon}m_{3/2}\xi(1-\rho)/\sqrt{2}$ . Again, using the central value of the oscillation parameters we obtain

$$\tilde{\epsilon} m_{3/2} \xi \sim 4 \times 10^{-10} \text{ GeV}. \quad (24)$$

This also results in a similar expression obtained in eq. (22). The model parameters in terms of the neutrino oscillation parameters in the absence of any CP violating Dirac phase becomes

$$\xi_i = \frac{\xi}{\sqrt{2}} \begin{pmatrix} 1.37 \\ 0.0264 \\ -0.349 \end{pmatrix}, \quad \tilde{\epsilon}_i = \frac{\tilde{\epsilon}}{\sqrt{2}} \begin{pmatrix} -0.27 \\ 0.922 \\ -1.036 \end{pmatrix}. \quad (25)$$

---

<sup>3</sup>In ref. [19] a similar form of the neutral fermion mass matrix was obtained but with different  $R$ -charge assignment. In that case, sneutrino VEV can be large and sneutrino can play the role of a down-type Higgs field. This feature is vastly different from our scenario where the sneutrino VEV is small because it is an  $R$ -breaking effect. Assorted phenomenology and collider implications of these two scenarios are also contrasting.



In this section we use the linear see-saw simplification where the  $(M_N)^{12}$  element is either much larger or smaller compared to the  $(M_N)^{13}$  element in eq. (20). The latter limit also solves a major drawback of prototypical MSSM-RPV scenario where  $\epsilon_i$  has to be very small in order to fit neutrino masses and mixing.

Apart from the tree level contribution there are one-loop corrections to the neutrino masses. There are two such dominant contributions from the  $BB$ - and the  $\mu B$ -loops [32], which in the decoupling regime ( $m_A \approx m_H$  and  $\sin^2(\alpha - \beta) \approx 1$ ) can be written in simplified forms as

$$\begin{aligned} (m_\nu)_{ij}^{BB} &\simeq \frac{g^2}{4(16\pi^2)^2} \left[ \frac{\cos^2(\alpha - \beta)}{\cos^2 \beta} \right] \left[ \frac{\tilde{\epsilon}_i \tilde{\epsilon}_j m_{3/2}^3}{\tilde{m}^2} \right], \\ (m_\nu)_{ij}^{\mu B} &\simeq \frac{g^2}{4(16\pi^2)^2} \left[ \frac{\cos(\alpha - \beta)}{\cos \beta} \right] \left[ \frac{\tilde{\epsilon}_i \tilde{\epsilon}_j m_{3/2}^2}{\tilde{m}} \right]. \end{aligned} \quad (26)$$

Thus even for  $\tilde{\epsilon}_i$ 's  $\sim \mathcal{O}(1)$  with gravitino mass around few hundred keV, these loop contributions turn out to be negligible, i.e.,  $m_\nu \sim \mathcal{O}(10^{-13})$  GeV.

#### 4. Gravitino as dark matter

In order to check the viability of gravitino DM, first we will discuss the production of gravitino and then its possible decay mode to adjudge whether the decay life-time of the gravitino is large enough to be a DM candidate or not. Then we will also discuss the constraints on photon flux emitting out of a gravitino decay.

##### 4.1. Production of gravitinos

In our scenario, gravitinos are produced in the thermal plasma by interacting with other SUSY particles, and their thermal relic density can be written as [33, 34]

$$\Omega_{3/2} h^2 \sim 0.1 \left( \frac{1 \text{ GeV}}{m_{3/2}} \right) \left( \frac{T_R}{10^7 \text{ GeV}} \right) \left( \frac{m_{\tilde{g}}}{2 \text{ TeV}} \right)^2, \quad (27)$$

where  $\Omega_{3/2} h^2 \sim 0.1199$  [2]. The thermal production of gravitinos would require  $T_R > m_{\tilde{g}} \sim 2 \text{ TeV}$ <sup>4</sup> otherwise the production of gravitinos would be exponentially (Boltzmann) suppressed [35]. Even though this outcome heavily depends on the exact SUSY spectrum, the constraint on the reheating temperature can be roughly satisfied with

$$m_{3/2} > 200 \text{ keV}. \quad (28)$$

Hence, the gravitino would constitute a cold DM [36]. If one requires more lighter gravitinos, the reheating temperature would naturally be lowered than the present consideration and thermal production of gravitinos becomes questionable. Nevertheless, it is always probable to have a late time entropy creation from the decay of other relics which can washout the gravitino density [37], and may relax the stringent constraint on the gravitino mass. Such low reheating temperature is also troublesome from the perspective of thermal leptogenesis [38]. Nevertheless, if the reheating temperature is above the electroweak phase transition, it can leave room for electroweak baryogenesis [39–42]. Another viable option for baryogenesis is the Affleck-Dine mechanism [43].

---

<sup>4</sup>In our framework, gluinos are very heavy compared to the electro-weak gauginos. Thus a pair of gluinos would decay to  $q\bar{q}\tilde{\chi}_1^0$ , where  $q$ 's are the first two generation quarks. For this kind of scenario, LHC provides stringent constraints on the gluino mass which is close to 2 TeV [6, 7].



#### 4.2. Gravitino decay width and life-time

Broken  $R$ -parity implies that the gravitino would most certainly decay. A light gravitino with mass  $\mathcal{O}(1 - 100)$  keV will dominantly decay to a neutrino and a photon. This decay width of gravitino is well approximated as [44]

$$\Gamma(\tilde{G} \rightarrow \nu_i \gamma) \simeq \frac{1}{32\pi} |U_{\gamma\nu}|^2 \frac{m_{3/2}^3}{M_P^2}, \quad (29)$$

where  $|U_{\gamma\nu}|^2 = \sum_{a=i+4} |\cos\theta_W Z_{a1} + \sin\theta_W Z_{a2}|^2$ , and  $M_P \sim 2.4 \times 10^{18}$  GeV is the reduced Planck mass. In a decoupled wino-higgsino scenario,  $|U_{\gamma\nu}|^2 = \cos^2\theta_W \sum_{a=i+4} |N_{a1}|^2$ , the physical light neutrino state can be approximated as

$$|\nu_m\rangle = \frac{|\nu_i\rangle}{\sqrt{1 + \left(\frac{\xi v_u}{M_D}\right)^2 + \left(\frac{g' v_i}{\sqrt{2} M_D}\right)^2}} + \frac{-\frac{g' \tilde{\epsilon}_i m_{3/2}}{\sqrt{2} \tilde{m}} |\tilde{b}\rangle}{\sqrt{1 + \left(\frac{\xi v_u}{M_D}\right)^2 + \left(\frac{g' v_i}{\sqrt{2} M_D}\right)^2}} + \frac{\frac{\xi_i v_u}{M_D} |\tilde{S}\rangle}{\sqrt{1 + \left(\frac{\xi v_u}{M_D}\right)^2 + \left(\frac{g' v_i}{\sqrt{2} M_D}\right)^2}}, \quad (30)$$

where  $m = 1, 2, 3$ . Using the above basis, the gravitino decay width given in eq. (29) can be further simplified as

$$\Gamma(\tilde{G} \rightarrow \nu_i \gamma) \simeq \frac{3}{32\pi} \left[ \frac{g' v_i}{\sqrt{2} M_D} \right]^2 \frac{m_{3/2}^3}{M_P^2} \cos^2\theta_W, \quad (31)$$

which in turn leads to the lifetime of gravitino decay

$$\tau_{3/2} \simeq \frac{5 \times 10^{19}}{\tilde{\epsilon}^2 m_{3/2}^5} \text{ sec}. \quad (32)$$

Now from this analysis it is obvious that this decay lifetime is much larger than the age of the Universe ( $4.32 \times 10^{17}$  sec) unless the gravitino is heavy ( $\geq \mathcal{O}(10)$  GeV) and  $\tilde{\epsilon} \sim \mathcal{O}(1)$ , unlike our choice, and this assures gravitino to be an excellent DM candidate.

#### 4.3. Photon flux from gravitino decays

As we have noted, the dominant decay mode of gravitino contains photon in final state, the constraints from the abundance of the photon production must be adjudged. This photon flux would most definitely show a peak at the energy  $E_\gamma = m_{3/2}/2$  and the maximum flux associated with it [44] can be deduced as

$$\begin{aligned} F_{\gamma, \text{ max}} &= E_\gamma \frac{dF}{dE_\gamma d\Omega} \bigg|_{E_\gamma = \frac{m_{3/2}}{2}} \\ &\simeq 2.89 \times 10^{-14} \tilde{\epsilon}_i^2 \left[ \frac{m_{3/2}}{200 \text{ keV}} \right]^2 \left[ \frac{\Omega_{3/2}}{0.3} \right] \left[ \frac{h}{0.7} \right] (\text{cm}^2 \cdot \text{str} \cdot \text{sec})^{-1}. \end{aligned} \quad (33)$$

Even for  $\tilde{\epsilon} \sim \mathcal{O}(1)$ , i.e.,  $\epsilon \sim \tilde{m}$ , which is the maximum value that we have accounted for, the flux is well within the bounds as analysed in [45].

### 5. Collider Phenomenology

In our model, as an artifact of  $R$ -symmetry breaking gravitino is the LSP. With suitable choice of parameters, we can choose a valid SUSY spectrum with lightest neutralino to be the next-to-minimal supersymmetric particle (NLSP).

Thus we would be left with a scenario where all the supersymmetric particles decay to the lightest neutralino which decays to a gravitino accompanied by either photon or  $Z$ -boson or Higgs. Since such interactions are suppressed by the Planck scale, the resulting decay width will be very small, i.e., corresponding

lifetime is too large to decay within the collider. In addition, the NLSP also undergoes  $R$ -parity violating decay modes primarily in the following channel

$$\tilde{\chi}_{1,2}^0 \rightarrow h \nu_i, \gamma \nu_i. \quad (34)$$

The dominant decay width is noted down as [46]

$$\Gamma(\tilde{\chi}_i^0 \rightarrow h\nu) = \frac{\alpha m_{\tilde{\chi}_i^0}}{16 \sin^2 \theta_W} \left| \left\{ \frac{\xi_i \cos \alpha}{\sqrt{2}} + \frac{g' \tilde{\epsilon} m_{3/2}}{\sqrt{2} \tilde{m}} \right\} N_{i3} N_{11} \right|^2 \left[ 1 - \frac{m_h^2}{m_{\tilde{\chi}_i^0}^2} \right]^2, \quad (35)$$

In our case, the dominant decay mode of this light neutralino is to a neutrino and a Higgs boson leading to an interesting di-Higgs signature at the colliders [47]. Charged lepton flavour violation constraints [48] put a limit on the parameter  $\xi \sim \mathcal{O}(10^{-2})$  or less [19]. This implies the decay length would be around 0.1 mm to a few mm. As an example, we show two representative points in the parameter space. (i)  $\tilde{\epsilon}, \xi \sim 10^{-3}$  and  $m_{3/2} \sim 100$  keV and (ii)  $\tilde{\epsilon} \sim 1$ ,  $m_{3/2} \sim 0.1$  GeV and  $\xi \sim 10^{-9}$ . Such different set of parameters provide distinct signatures at the colliders. The later consideration also solves the ‘small  $\epsilon$ ’ issue in traditional bi-linear “RPV” scenarios. A measurable decay gap is a distinctive feature of this scenario. In the wino-higgsino decoupled scenario the lightest neutralino is primarily an admixture of the bino-singlino states and hence the decay modes to  $Z\nu$  and  $W^\pm \ell^\mp$  which are highly suppressed. One can easily have light higgsinos in this framework without much modifications in the neutrino sector and would also have interesting collider implications [49].

In presence of small  $R$ -breaking parameter, a small mass splitting between the neutralinos and sneutrinos are induced due to the pseudo-Dirac nature of neutralinos and mixing of sneutrinos with neutral Higgs bosons [50] respectively. Thus we can have neutralino oscillations at the LHC [51] which can be identified in the distribution of the displaced vertex lengths. There is also a possibility of having induced sneutrino-antisneutrino oscillations [52] which are an excellent probe for the lepton number violation. The bi-linear “RPV” violating but  $R$ -symmetry preserving term in the superpotential with the sneutrino VEV would also lead to trilinear couplings involving lepton (quark) and slepton (squark) fields which are also an interesting channel to look for [53]. All these together make our scenario phenomenologically rich and explorable at the LHC.

## 6. Conclusion

Neutrinos have always been the elusive key to understand the physics beyond the Standard Model. We believe the lack of our definiteness regarding neutrino mass generation mechanism is one of the hurdle to propel the search for BSM physics. Keeping trust on theoretical construction and motivations of supersymmetry, in this paper we come up with an  $R$ -symmetric SUSY scenario where gauginos are Dirac type and there is a natural emergence of bi-linear “ $R$ -parity violation” as well. While constructing the superpotential within our framework we note that there are  $R$ -symmetry conserving terms which are interestingly “ $R$ -parity violating”. This is indeed an interesting outcome of our scenario. Our prime aim in this paper is to explain how active light neutrino masses and mixing can be generated. In the process we also discuss one of the important issue: the generation of a Higgs mass around 125 GeV. Then we motivate the requirement of  $R$ -symmetry breaking and that also through the conformal compensator field. These  $R$ -symmetry breaking terms are instrumental to generate tiny VEV of sneutrinos which is suitable for neutrino mass and mixing generation. Though we have the full neutral fermion mass matrix, but for the sake of neutrino mass generation we work with wino-higgsino decoupling limit without loss of generality. Then we deduce the constraints on relevant superpotential parameters for both normal and inverted hierarchies. After successful generation of neutrino masses using tree level operators, we also show that the loop induced contributions are small for our choice of parameter space and can be ignored. In our scenario gravitino is the LSP and a possible dark matter candidate. We explore this possibility by considering the production and decays of gravitino. These lead to the constraints on masses and some couplings relevant for neutrino masses as well. We then also take care of the photon flux constraint. It is worthy to mention that all these constraints are compatible

among themselves leading to a successful scenario we are looking for. At the end we conclude by providing some directions along which this scenario can be looked for at the colliders, like LHC.

## 7. Appendix

Here, we discuss the limit when  $m_D \ll m_{\tilde{f}_0}$  and hence Type-I see-saw approximation can be used. We again work under the simplified assumption when the wino is heavy and effectively decoupled. In that scenario the neutral fermion mass matrix in the basis  $(\nu_i, \tilde{b}, \tilde{S}, \tilde{H}_u^0, \tilde{R}_d^0, \tilde{H}_d^0, \tilde{R}_u^0)$  can be obtained from the neutral fermion mass matrix  $M_N$ . Using the see-saw relation we find the neutrino mass-matrix elements can be written in a closed form as

$$(m_\nu)_{ij} = \tilde{a} \epsilon_i \epsilon_j + \tilde{b} [\epsilon_i \xi_j + \epsilon_j \xi_i] + \tilde{c} \xi_i \xi_j, \quad (36)$$

where

$$\begin{aligned} \tilde{a} &= \frac{m_{3/2}}{\tilde{m}^2} [g'^2 \lambda^2 - \sqrt{2} g' \lambda], \\ \tilde{b} &= \frac{m_{3/2}}{2\tilde{m}} [\sqrt{2} g' - g'^2 \lambda], \\ \tilde{c} &= -\frac{m_{3/2}}{16\pi^2}. \end{aligned} \quad (37)$$

Here,  $\tilde{m}$  is the generic  $R$ -conserving mass scale. Eq. (36) also represents a striking feature, as in usual bi-linear RPV models such a structure only arises at the one-loop level [54–58]. Moreover, in order to fit the neutrino oscillation data one requires unnatural fine-tuning among parameters. However, in this scenario not only all the neutrinos acquire mass at the tree level but the smallness of the neutrino mass can be entangled to the gravitino mass itself. Similarly, when both the wino and the higgsinos are decoupled one obtains the effective neutrino mass matrix as

$$(m_\nu)_{ij} = \tilde{b}' [\epsilon_i \xi_j + \epsilon_j \xi_i] + \tilde{c} \xi_i \xi_j, \quad (38)$$

with

$$\tilde{b}' = \frac{g' m_{3/2}}{\sqrt{2} \tilde{m}}. \quad (39)$$

## Acknowledgements

It is a pleasure to thank Sourov Roy for many helpful discussions and insights. SC would also like to thank the hospitality of IIT Kanpur and HRI Allahabad where a part of the project was completed.

## References

- [1] I. Esteban, M. C. Gonzalez-Garcia, M. Maltoni, I. Martinez-Soler and T. Schwetz, “Updated fit to three neutrino mixing: exploring the accelerator-reactor complementarity,” arXiv:1611.01514 [hep-ph].
- [2] P. A. R. Ade *et al.* [Planck Collaboration], “Planck 2013 results. XVI. Cosmological parameters,” *Astron. Astrophys.* **571**, A16 (2014),
- [3] A. C. M. Correia *et al.*, “The HARPS search for southern extra-solar planets XIX. Characterization and dynamics of the GJ876 planetary system,” *Astron. Astrophys.* **511**, A21 (2010)
- [4] R. Barbier *et al.*, “R-parity violating supersymmetry,” *Phys. Rept.* **420**, 1 (2005).
- [5] J. L. Feng, “Dark Matter Candidates from Particle Physics and Methods of Detection,” *Ann. Rev. Astron. Astrophys.* **48**, 495 (2010).
- [6] The ATLAS collaboration [ATLAS Collaboration], “Further searches for squarks and gluinos in final states with jets and missing transverse momentum at  $\sqrt{s}=13$  TeV with the ATLAS detector,” ATLAS-CONF-2016-078.
- [7] CMS Collaboration [CMS Collaboration], “Search for supersymmetry in events with jets and missing transverse momentum in proton-proton collisions at 13 TeV,” CMS-PAS-SUS-16-014.

- [8] P. J. Fox, A. E. Nelson and N. Weiner, “Dirac gaugino masses and supersoft supersymmetry breaking,” JHEP **0208**, 035 (2002).
- [9] G. D. Kribs and A. Martin, “Supersoft Supersymmetry is Super-Safe,” Phys. Rev. D **85**, 115014 (2012).
- [10] G. D. Kribs, E. Poppitz and N. Weiner, “Flavor in supersymmetry with an extended R-symmetry,” Phys. Rev. D **78**, 055010 (2008).
- [11] C. Frugiuele and T. Gregoire, “Making the Sneutrino a Higgs with a  $U(1)_R$  Lepton Number,” Phys. Rev. D **85**, 015016 (2012).
- [12] E. Bertuzzo and C. Frugiuele, “Fitting Neutrino Physics with a  $U(1)_R$  Lepton Number,” JHEP **1205**, 100 (2012).
- [13] C. Frugiuele, T. Gregoire, P. Kumar and E. Ponton, “ $L=R' - U(1)_R$  as the Origin of Leptonic ‘RPV’,” JHEP **1303**, 156 (2013).
- [14] C. Frugiuele, T. Gregoire, P. Kumar and E. Ponton, “ $L=R' - U(1)_R$  Lepton Number at the LHC,” JHEP **1305**, 012 (2013).
- [15] S. Chakraborty and S. Roy, “Higgs boson mass, neutrino masses and mixing and keV dark matter in an  $U(1)_R$  – lepton number model,” JHEP **1401**, 101 (2014).
- [16] S. Chakraborty, D. K. Ghosh and S. Roy, “7 keV Sterile neutrino dark matter in  $U(1)_R$  – lepton number model,” JHEP **1410**, 146 (2014).
- [17] S. Chakraborty, A. Datta and S. Roy, “ $h\gamma\gamma$  in  $U(1)_R$  -lepton number model with a right-handed neutrino,” JHEP **1502**, 124 (2015), Erratum: [JHEP **1509**, 077 (2015)].
- [18] S. Chakraborty, A. Datta, K. Huitu, S. Roy and H. Waltari, “Light top squarks in  $U(1)_R$ -lepton number model with a right handed neutrino and the LHC,” Phys. Rev. D **93**, no. 7, 075005 (2016).
- [19] P. Coloma and S. Ipek, “Neutrino masses from a pseudo-Dirac Bino,” Phys. Rev. Lett. **117**, no. 11, 111803 (2016).
- [20] R. Fok and G. D. Kribs, “ $\mu$  to  $e$  in R-symmetric Supersymmetry,” Phys. Rev. D **82**, 035010 (2010).
- [21] S. Y. Choi, D. Choudhury, A. Freitas, J. Kalinowski, J. M. Kim and P. M. Zerwas, “Dirac Neutralinos and Electroweak Scalar Bosons of  $N = 1/N = 2$  Hybrid Supersymmetry at Colliders,” JHEP **1008**, 025 (2010).
- [22] R. Fok, G. D. Kribs, A. Martin and Y. Tsai, “Electroweak Baryogenesis in R-symmetric Supersymmetry,” Phys. Rev. D **87**, no. 5, 055018 (2013).
- [23] S. Chakraborty, A. Chakraborty and S. Raychaudhuri, “Diphoton resonance at 750 GeV in the broken R-symmetric MSSM,” Phys. Rev. D **94**, no. 3, 035014 (2016).
- [24] K. Benakli, M. D. Goodsell and F. Staub, “Dirac Gauginos and the 125 GeV Higgs,” JHEP **1306**, 073 (2013).
- [25] M. D. Goodsell, M. E. Krauss, T. Mller, W. Porod and F. Staub, “Dark matter scenarios in a constrained model with Dirac gauginos,” JHEP **1510**, 132 (2015).
- [26] A. E. Nelson and T. S. Roy, “New Supersoft Supersymmetry Breaking Operators and a Solution to the  $\mu$  Problem,” Phys. Rev. Lett. **114**, 201802 (2015).
- [27] M. D. Goodsell, “Two-loop RGEs with Dirac gaugino masses,” JHEP **1301**, 066 (2013).
- [28] Z. Chacko, M. A. Luty, I. Maksymyk and E. Ponton, “Realistic anomaly mediated supersymmetry breaking,” JHEP **0004**, 001 (2000).
- [29] P. Diessner, J. Kalinowski, W. Kotlarski and D. Stöckinger, “Higgs boson mass and electroweak observables in the MRSSM,” JHEP **1412**, 124 (2014).
- [30] P. Diessner, J. Kalinowski, W. Kotlarski and D. Stöckinger, “Two-loop correction to the Higgs boson mass in the MRSSM,” Adv. High Energy Phys. **2015**, 760729 (2015).
- [31] M. B. Gavela, T. Hambye, D. Hernandez and P. Hernandez, “Minimal Flavour Seesaw Models,” JHEP **0909**, 038 (2009).
- [32] Y. Grossman and S. Rakshit, “Neutrino masses in R-parity violating supersymmetric models,” Phys. Rev. D **69**, 093002 (2004).
- [33] M. Bolz, A. Brandenburg and W. Buchmüller, “Thermal production of gravitinos,” Nucl. Phys. B **606**, 518 (2001) Erratum: [Nucl. Phys. B **790**, 336 (2008)].
- [34] J. Pradler and F. D. Steffen, “Thermal gravitino production and collider tests of leptogenesis,” Phys. Rev. D **75**, 023509 (2007).
- [35] N.-E. Bomark and L. Roszkowski, “3.5 keV x-ray line from decaying gravitino dark matter,” Phys. Rev. D **90**, 011701 (2014).
- [36] M. Fujii, M. Ibe and T. Yanagida, “Upper bound on gluino mass from thermal leptogenesis,” Phys. Lett. B **579**, 6 (2004).
- [37] E. A. Baltz and H. Murayama, “Gravitino warm dark matter with entropy production,” JHEP **0305**, 067 (2003).
- [38] M. Fujii and T. Yanagida, “Natural gravitino dark matter and thermal leptogenesis in gauge mediated supersymmetry breaking models,” Phys. Lett. B **549**, 273 (2002).
- [39] W. Fischler, G. F. Giudice, R. G. Leigh and S. Paban, “Constraints on the baryogenesis scale from neutrino masses,” Phys. Lett. B **258**, 45 (1991).
- [40] B. A. Campbell, S. Davidson, J. R. Ellis and K. A. Olive, “Cosmological baryon asymmetry constraints on extensions of the standard model,” Phys. Lett. B **256**, 484 (1991).
- [41] H. K. Dreiner and G. G. Ross, “Sphaleron erasure of primordial baryogenesis,” Nucl. Phys. B **410**, 188 (1993).
- [42] V. A. Kuzmin, V. A. Rubakov and M. E. Shaposhnikov, “On the Anomalous Electroweak Baryon Number Nonconservation in the Early Universe,” Phys. Lett. **155B**, 36 (1985).
- [43] I. Affleck and M. Dine, “A New Mechanism for Baryogenesis,” Nucl. Phys. B **249**, 361 (1985).
- [44] F. Takayama and M. Yamaguchi, “Gravitino dark matter without R-parity,” Phys. Lett. B **485**, 388 (2000).
- [45] P. Sreekumar *et al.* [EGRET Collaboration], “EGRET observations of the extragalactic gamma-ray emission,” Astrophys. J. **494**, 523 (1998).
- [46] S. Bobrovskiy, J. Hajer and S. Rydbeck, “Long-lived higgsinos as probes of gravitino dark matter at the LHC,” JHEP

- 1302**, 133 (2013).
- [47] S. Biswas, E. J. Chun and P. Sharma, “Di-Higgs signatures from R-parity violating supersymmetry as the origin of neutrino mass,” JHEP **1612**, 062 (2016).
  - [48] L. Bartoszek *et al.* [Mu2e Collaboration], “Mu2e Technical Design Report,” arXiv:1501.05241 [physics.ins-det].
  - [49] B. Mukhopadhyaya, S. Roy and F. Vissani, “Correlation between neutrino oscillations and collider signals of supersymmetry in an R-parity violating model,” Phys. Lett. B **443**, 191 (1998).
  - [50] E. J. Chun, “CP violation, sneutrino oscillation and neutrino masses in R-parity violating supersymmetric standard model,” Phys. Lett. B **525**, 114 (2002).
  - [51] Y. Grossman, B. Shakya and Y. Tsai, “Neutralino Oscillations at the LHC,” Phys. Rev. D **88**, no. 3, 035026 (2013).
  - [52] Y. Grossman and H. E. Haber, “Sneutrino mixing phenomena,” Phys. Rev. Lett. **78**, 3438 (1997).
  - [53] S. Roy and B. Mukhopadhyaya, “Some implications of a supersymmetric model with R-parity breaking bilinear interactions,” Phys. Rev. D **55**, 7020 (1997).
  - [54] M. Hirsch, M. A. Diaz, W. Porod, J. C. Romao and J. W. F. Valle, “Neutrino masses and mixings from supersymmetry with bilinear R parity violation: A Theory for solar and atmospheric neutrino oscillations,” Phys. Rev. D **62**, 113008 (2000) Erratum: [Phys. Rev. D **65**, 119901 (2002)].
  - [55] G. G. Ross and J. W. F. Valle, “Supersymmetric Models Without R-Parity,” Phys. Lett. **151B**, 375 (1985).
  - [56] J. R. Ellis, G. Gelmini, C. Jarlskog, G. G. Ross and J. W. F. Valle, “Phenomenology of Supersymmetry with Broken R-Parity,” Phys. Lett. **150B**, 142 (1985).
  - [57] M. A. Diaz, M. Hirsch, W. Porod, J. C. Romao and J. W. F. Valle, “Solar neutrino masses and mixing from bilinear R parity broken supersymmetry: Analytical versus numerical results,” Phys. Rev. D **68**, 013009 (2003) Erratum: [Phys. Rev. D **71**, 059904 (2005)].
  - [58] F. Bazzocchi, S. Morisi, E. Peinado, J. W. F. Valle and A. Vicente, “Bilinear R-parity violation with flavor symmetry,” JHEP **1301**, 033 (2013).